CERN-TH/99-84
IEM-FT-178/98
IFT-UAM/CSIC-99-11

F-term inflation in Superstring Theories

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Abstract

A supersymmetric inflationary stage dominated by an F-term has the problem that the flatness of the potential is spoiled by supergravity corrections, that is the slow-roll parameter η gets contributions of order unity. We show that in F-term inflationary models based on strings there is natural way of obtaining small values of η. This happens in models of hybrid inflation based on orbifold constructions, in which a modulus T field is responsible for the large value of the potential during inflation, and a second field ϕ with appropriate modular weight is responsible for the roll-over. We illustrate the mechanism with a model in which the inflaton potential is provided by gaugino condensation, leading to succesful inflation.

March, 1999

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The flatness and the horizon problems of the standard big bang cosmology are elegantly solved if during the evolution of the early Universe the energy density happens to be dominated by some form of vacuum energy and comoving scales grow quasi-exponentially [1]. An inflationary stage is also required to dilute any undesirable topological defects left as remnants after some phase transition taking place at early epochs. The vacuum energy driving inflation is generally assumed to be associated to some scalar field ϕ, the inflaton, which is initially displaced from the minimum of its potential. As a by-product, quantum fluctuations of the inflaton field may be the seeds for the generation of structure. The levels of density and temperature fluctuations observed in the present Universe, $\delta\rho/\rho \sim 10^{-5}$, require the inflaton potential to be extremely flat. Actually, for inflation to be viable, there are two parameters which must be small with respect to unity to insure sufficient slow roll in the inflationary potential V. These are the so called $\epsilon = \frac{1}{2}M_{\rm Pl}^2(V'/V)^2$ and $\eta = M_{\rm Pl}^2 V''/V$ parameters. Here $M_{\rm Pl} = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. The value of η in particular is constrained by the observational bound in the spectral index $|n-1| \lesssim 0.2$. Since in an inflationary scenario $n = 1 - 6\epsilon + 2\eta$, this requires $|\eta| \lesssim 0.1$ (the contribution proportional to ϵ is usually negligible). Note that η is essentially given by the effective inflaton mass during inflation.

The required extreme flatness of the inflaton potential is the main reason why inflation is more natural in the context of supersymmetric (SUSY) theories. Flat directions are frequent in supersymmetric theories, and they are stable under radiative corrections, as long as SUSY is not broken. However, the previous argument does not hold if SUSY is broken, which necessarily occurs during inflation, when there is a vacuum energy density $\langle V\rangle_{\rm in} \sim H^2 M_{\rm Pl}^2 > 0$. Thus, one generically expects effective soft terms, in particular soft scalar masses, of order H, spoiling the desired flatness of the potential.

The supergravity potential V consists of two pieces — the so-called D-term and F-term. The latter is given by

$$V_F = e^G \left(G_{\bar{j}} K^{\bar{j}l} G_l - 3M_{\rm Pl}^4 \right) = F^{\bar{l}} K_{j\bar{l}} F^j - 3e^G M_{\rm Pl}^4 \ . \tag{1}$$

Here $G = K M_{\rm Pl}^{-2} + \ln(|W|^2 M_{\rm Pl}^{-6})$, where W is the superpotential, $K^{\bar{l}j}$ is the inverse of the Kähler metric $K_{j\bar{l}} \equiv \partial K/\partial\phi_j \partial\bar{\phi}_l$, ϕ_j are the (scalar components) of the chiral superfields and $F^j = e^{G/2} K^{\bar{k}j} G_{\bar{k}} M_{\rm Pl}^3$ are the corresponding auxiliary fields. In the following, we will assume that during inflation the potential is dominated by the F-term, i.e. $V \simeq V_F$, which is the so-called F–inflation. Notice that the F-term depends only upon two objects — the

superpotential W, an holomorphic function of the scalar fields, and the Kähler potential K, a real function. During inflation

$$\langle V\rangle_{\rm in} \equiv V_0 = 3H^2 M_{\rm Pl}^2, \tag{2}$$

which implies that some F fields are different form zero, thus breaking supersymmetry.

It is easy to check that this breaking of SUSY generically gives a sizeable soft mass-squared to the inflaton. Using a basis where the fields have canonical normalization at the origin ($K_{i\bar{j}} = \delta_{ij}$ at $\phi_i = 0$), a generic scalar mass $m_i^2 = \partial^2 V/\partial\phi_i\partial\phi_i^*$ is given by

$$m_i^2 = M_{\rm Pl}^{-2} V_0 + e^K \sum_{nm} K_{ii^*}^{n^*m} W_n W_m^* - M_{\rm Pl}^{-2} e^K |W_i|^2 + e^K \sum_n |W_{ni}|^2 + \cdots . \tag{3}$$

The contribution of the first term to $\eta \equiv M_{\rm Pl}^{-2} m_\phi^2 V^{-1}$ is precisely 1. For generic models, the contribution of the second term to η is also $\mathcal{O}(1)$. So, even if in the global SUSY limit the inflaton is not responsible for the SUSY breaking ($W_\phi = 0$), remaining as a flat direction during inflation ($W_{n\phi} = 0$), we see from eq.(3) that the mass-squared of the inflaton is expected to be $\mathcal{O}(M_{\rm Pl}^{-2} V)$, and thus $\eta = \mathcal{O}(1)$. Therefore, it is generally difficult to naturally implement a slow-roll inflation in the context of supergravity (recall we need $|\eta| \lesssim 0.1$).

Solutions to the η-problem already exist in the literature [4]. Among them, D-term inflation seems to be a promising one [5]. In this paper, however, we show that in F-term inflationary models based on (weakly coupled) strings there is a natural way of obtaining small values of η. This happens in models of hybrid inflation [6] based on orbifold constructions, in which a modulus T field is responsible for the large value of the potential V during inflation, and a second field ϕ with modular weight $n_\phi = -3$ is responsible for the rollover [7].

String models of F-term inflation scenarios present special characteristics because the Kähler potential, K, which plays a crucial role, is greatly constrained. In orbifold constructions, the tree–level Kähler potential is given by [8]

$$K = -\log(S+\bar{S}) - 3\log(T+\bar{T}) + \sum_j (T+\bar{T})^{n_j} |\phi_j|^2 + \mathcal{O}\left(\frac{|\phi_j|^4}{M_{\rm Pl}^{-2}}\right)) \ . \tag{4}$$

Here S is the dilaton and T denotes generically the moduli fields, all in units of $M_{\rm Pl}$, ϕ_j are the chiral fields and n_j the corresponding modular weights. The latter depend on the type of orbifold considered and the twisted sector to which the field belongs. The possible values of n_j for Z_N orbifolds are $n_j = -1, -2, -3, -4, -5$. The discrete character of n_j will

play a relevant role later on. Eq.(4) is written with the usual simplification of considering a single "overall modulus" T. $\langle S\rangle$ and $\langle T\rangle$ are expected to have determined values of $\mathcal{O}(1)$ in Planck units at low energy. This comes from the fact that $\langle S\rangle$ and $\langle T\rangle$ have precise physical meanings. Namely, $\langle S\rangle$ is the value of the unified gauge coupling constant and $\langle T\rangle$ is the squared radius of the compactified space, both in Planck units.

From eq.(4) it is easy to understand that T is a good candidate for providing the vacuum energy during inflation (i.e. $F^T \neq 0$). The reason is the coupling between ϕ and T in the quadratic term ($\propto |\phi|^2$) of K. Notice that, without this coupling, if ϕ corresponds to a flat direction in global SUSY breaking ($W_\phi = W_{n\phi} = 0$), then $m_\phi^2 = V_0 M_{\rm Pl}^{-2}$ (see eq.(3)) and thus $\eta = 1$.

If SUSY is broken by the T (and/or S) fields, the corresponding soft terms, in particular mass terms, for ϕ are straightforwardly extracted from Eqs.(1) and (4). This was precisely the sort of scenario considered in Refs. [9,10]. Although the motivation of these works was different, namely to study the form of the soft breaking terms at low–energy with $m_{3/2} = \mathcal{O}(1\ \text{TeV})$, their results are applicable here. The only difference here is that the scale of the breaking is much higher and the non-vanishing cosmological constant, $V_0 > 0$, plays a major role. In particular, the effective ϕ mass squared, m_ϕ^2, of a chiral field ϕ with modular weight n_ϕ, for the canonically normalized field $(K_{\phi\bar\phi})^{1/2}\phi$, is given by [10]

$$m_\phi^2 = m_{3/2}^2\left\{(3 + n_\phi \cos^2\theta)C^2 - 2\right\}\ , \tag{5}$$

where the effective gravitino mass during inflation, $m_{3/2}$, is defined by

$$m_{3/2}^2 = e^G M_{\rm Pl}^2 = e^{K/M_{\rm Pl}^2}|W|^2 M_{\rm Pl}^{-4}\ . \tag{6}$$

(since $V_0 = F^{\bar l}K_{j\bar l}F^j - 3m_{3/2}^2 M_{\rm Pl}^2$, unless there is some fine–tuning, $m_{3/2}^2$ is at most $\mathcal{O}(V_0)M_{\rm Pl}^{-2}$); $\tan^2\theta = (K_{S\bar S}/K_{T\bar T})\left|F^S/F^T\right|^2$ and $C^2 = 1 + [V_0/(3M_{\rm Pl}^2 m_{3/2}^2)]$. Notice that $C^2 > 1$.

From (5) we conclude the following. If $\cos^2\theta = 0$, the case of S–driven inflation, then $m_\phi^2 = m_{3/2}^2 + V_0 M_{\rm Pl}^{-2}$. Hence, $\eta = \mathcal{O}(1)$ as expected from previous arguments. Let us examine now the more promising case of T-driven inflation, where $\cos^2\theta = 1$. Note that in this case $V_0 = K_{T\bar T}|F^T|^2 - 3m_{3/2}^2 M_{\rm Pl}^2 = 3H^2 M_{\rm Pl}^2$. If it happens that $K_{T\bar T}|F^T|^2 = \mathcal{O}(m_{3/2}^2)M_{\rm Pl}^2$, what means $C^2 = \mathcal{O}(1)$, then still one obtains the usual result of $m_\phi^2 = \mathcal{O}(m_{3/2}^2) = \mathcal{O}(H^2)$, and thus $\eta = \mathcal{O}(1)$. However, it may perfectly happen that

$$K_{T\bar T}|F^T|^2 \gg m_{3/2}^2 M_{\rm Pl}^2. \tag{7}$$

Then $V_0 \gg \mathcal{O}(m_{3/2}^2 M_{\rm Pl}^2)$ and $C^2 \gg 1$. Then Eq. (5) simplifies to

$$m_\phi^2 \simeq \frac{1}{3} V_0 M_{\rm Pl}^{-2} (3 + n_\phi) = H^2 (3 + n_\phi). \tag{8}$$

Hence, for $n_\phi = -3$ we get $m_\phi^2 \ll \mathcal{O}(H^2)$ and thus $\eta \ll 1$, as desired†.

Therefore, we have shown that the possibility of a very small mass m_ϕ^2 can occur in T driven F-term inflation, if $V_0 \gg \mathcal{O}(m_{3/2}^2 M_{\rm Pl}^2)$ (namely $C^2 \gg 1$) and the modular weight ϕ is $n_\phi = -3$. Notice that this possibility appears with no fine tunning at all, thanks to the discrete character of n_ϕ, thus avoiding the slow-roll η-problem. In this scenario the energy density of the vacuum is provided by the modulus T and the slow-roll field is the ϕ field. We are envisaging, therefore, an hybrid inflationary scenario where the dynamics of the system will make the energy density of the vacuum disappear for some critical value of the field ϕ.

To get these nice results we have used the form of the Kähler potential of eq.(4), which contains two main approximations: *(i)* it is a tree-level expression, and *(ii)* it is leading order in the expansion in powers of $|\phi|^2$.

Concerning approximation *(i)*, perturbative corrections to eq.(4) are known at one-loop level [8] and are small, so they do not affect any of the results presented here. More precisely, at one loop the correction consists in replacing in Eq. (4)

$$S + \bar{S} \rightarrow Y = S + \bar{S} - \frac{\delta_{GS}}{8\pi^2} \ln(T + \bar{T}), \tag{9}$$

where δ_{GS} is the Green-Schwarz parameter, which is negative and $|\delta_{GS}| \leq \mathcal{O}(10)$. With this replacement, Eq. (5) becomes

$$m_\phi^2 = m_{3/2}^2 \left\{ \left[3 + n_\phi \left(1 - \frac{\delta_{GS}}{24\pi^2 Y} \right)^{-1} \cos^2\theta \right] C^2 - 2 \right\} . \tag{10}$$

In the case we identified above, in which $\cos^2\theta = 1$, $C^2 \gg 1$ and $n_\phi = -3$, from (10) we obtain

$$m_\phi^2 \simeq -H^2 \frac{\delta_{GS}}{8\pi^2 Y} \simeq -H^2 \frac{\delta_{GS}}{32\pi^2}, \tag{11}$$

where we have used that Y is the inverse of the unified coupling constant g, namely $Y \simeq 2/g^2$, so we expect $\langle Y \rangle \simeq 4$. Eq.(11) shows that $m_\phi^2 < 0.1\ H^2$, which leads to values

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†To get this result we have assumed eq.(7), which is equivalent to $|W_T| \gg \frac{9}{(T+\bar{T})^2}|W|^2$. This is easily fulfilled for slightly large values of T. E.g. if $W \propto [\eta(T)]^{-6}$, as dictated by target-space modular invariance, then $|W_T| \gg \frac{5}{2}|W|^2$ and $|\eta| < 0.1$ for Re$T > 2.5$.

of η small enough to allow for the implementation of inflation. On the other hand, non-perturbative corrections are very poorly known (see e.g. Ref. [11] for an analysis of their possible phenomenological significance).

Concerning approximation *(ii)*, the Kähler potential will generically contain terms of the form

$$K = K_0 + K_{\phi\bar{\phi}}|\phi|^2 + \lambda_4|\phi|^4 + \lambda_6|\phi|^6 + ..., \tag{12}$$

where the first two terms are the ones explicited in eq.(4), and one expects $\lambda_n = \mathcal{O}(M_{\rm Pl}^{2-n})$. These extra terms modify the potential to be

$$V = V_0 \left(1 + \kappa_1^4|\phi|^4 + \kappa_2^6|\phi|^6 + ...\right) \ , \tag{13}$$

where the coefficients κ_i are all of the order of $M_{\rm Pl}^{-1}$. For $\phi \lesssim 0.1 M_{\rm Pl}$ the $|\phi|^4, |\phi|^6, ...$ terms become negligible (this is a common situation and, as we will see, it is the case of the particular model we present below). For example, keeping only the $|\phi|^4$ term above, and ignoring the phase of ϕ, the ϵ and η parameters are given by

$$\epsilon = \frac{8M_{\rm Pl}^2\kappa_1^8\phi^6}{(1+\kappa^4\phi^4)^2} \ , \quad \eta = \frac{12M_{\rm Pl}^2\kappa_1^2\phi^2}{1+\kappa_1^2\phi^4}. \tag{14}$$

Then, if $\phi \lesssim 0.1 M_{\rm Pl}$ we get $\epsilon \simeq 8\kappa_1^6\phi^6 \lesssim 10^{-5}$, $\eta \simeq 12(\phi/M_{\rm Pl})^2 \lesssim 0.1$.

It is amusing to note that the exact (tree-level) form of $K(\phi, \bar{\phi})$ is not known, but has been conjectured [12]. For our case of interest, $n_\phi = -3$, it simply becomes

$$K = -\log(S+\bar{S}) - \log\left[(T+\bar{T})^3 - |\phi|^2\right]. \tag{15}$$

Then, assuming $V_0 \gg M_{\rm Pl}^2 m_{3/2}^2$ (i.e. $C^2 \gg 1$), it turns out that all the $|\phi|^4, |\phi|^6, ...$ contributions in (13) are exactly vanishing!

We have discussed a natural way to solve the slow-roll problem (i.e. to get $|\eta| \ll 1$ during inflation). However, without knowing the T-dependent potential we cannot know the value of the field ϕ at the end of inflation ϕ_e. Neither we can calculate its value 60 inflation-folds before the end, ϕ_{60}. This is important to get more definite results. Next we show a model based on gaugino condensation in which the previous mechanism to get $|\eta| \ll 1$ is naturally implemented, allowing for further predictions.

Suppose the hidden sector of the theory has an $SU(N_c)$ gauge group with N_f flavours. Then the gaugino condensation effects generate a non-perturbative superpotential W_{np} of the form [13]

$$W_{np} = \det \mathcal{M}^{1/N_c}\ f(T). \tag{16}$$

Here $\mathcal{M}$ is the effective "quark" mass matrix and $f(T) \sim e^{-8\pi^2/N_c g^2}[\eta(T)]^\alpha$, where $\alpha = \mathcal{O}(1)$. Suppose that $\mathcal{M}$ has an invariant contribution and a field-dependent contribution, i.e.

$$\mathcal{M} = M - \lambda A^q\ , \tag{17}$$

where M is a constant mass term (typically $M = \mathcal{O}(M_{\rm Pl})$), A is a scalar field, λ is a coupling constant and q is a numerical exponent ($q = 1$ or $q = 2$ is okay for our purposes). The A field can also have perturbative interactions with the inflaton ϕ (i.e. a field with $n_\phi = -3$), in particular $\sim A^2\phi^r$, where r is an undetermined exponent. Then the relevant superpotential gets the form

$$W_{np} = [M - \lambda A^q]^{N_f/N_c}\ f(T)\ +\ \lambda' A^2 \phi^r\ , \tag{18}$$

where λ' is a coupling constant. As we will see, we need $r \geq 2$ for our model to work.

It is easy to check that the supergravity potential, V_{SUGRA}, derived from eq.(18) has a *global minimum* at

$$\lambda A^q = M \qquad \phi = 0\ . \tag{19}$$

Notice that at this point $\det \mathcal{M} = 0$ and $V = 0$, and supersymmetry is unbroken since

$$W = W_T = W_A = W_\phi = 0\ . \tag{20}$$

On the other hand, for large $\langle\phi\rangle$ the A field gets an effective large mass and, hence, the *effective minimum* of V_{SUGRA} is at

$$A \simeq 0\ . \tag{21}$$

So, the ϕ potential in this region is simply given by

$$V(\phi) = K_{T\bar{T}} |F^T|^2 \equiv V_0\ , \tag{22}$$

which is flat, leading to an inflationary process. Actually, the potential is not exactly flat. As explained above, perturbative contributions to the Kähler potential give to ϕ a positive tiny mass $m_\phi^2 \simeq H^2 |\frac{\delta_{GS}|}{32\pi^2}$ (see eqs.(9–11)). So, the inflaton ϕ rolls slowly towards the origin. Inflation ends when the effective A mass becomes unimportant. This occurs for a value of $\phi = \phi_e$ given by

$$\phi_e \simeq V_0^{1/2r} {M_{\rm Pl}}^{1-2/r} \ . \tag{23}$$

For $\phi < \phi_e$, ϕ and A move quickly towards the true (global) minimum and inflation ends. The value of ϕ 60 inflation-folds before the end, ϕ_{60}, is given by

$$\phi_{60} \simeq e^{|\delta_{GS}|/16} \phi_e \simeq 2\phi_e \ . \tag{24}$$

This is crucial to evaluate the primordial spectrum of scalar perturbations

$$\mathcal{P}_\xi = \frac{\kappa^6}{3(2\pi)^2} \frac{V_{60}^3}{(V_{60}')^2} \ . \tag{25}$$

According to observations $\mathcal{P}_\xi \simeq (5 \times 10^{-5})^2$, requiring

$$V_0^{1/4} = \left[2 \times 10^{-5}\right]^{\frac{r}{2(r-1)}} M_{\rm Pl} \ . \tag{26}$$

More precisely, for $r \geq 2$ this translates into

$$5 \times 10^{13} {\rm GeV} \leq V_0^{1/4} \leq 10^{16} {\rm GeV} \ , \tag{27}$$

where the lower bound corresponds to the simplest case $r = 2$. We finally note that these numbers for V_0 are perfectly natural in our model. More precisely, V_0 is given by eqs.(22,18,21), so typically is $V_0^{1/4} \sim e^{-4\pi^2/N_c g^2} \mathcal{O}(M_{\rm Pl})$. Thus, the suppression needed to bring $V_0^{1/4}$ into the range of eq.(27) is provided by $N_c \sim 8$, a perfectly reasonable number. It is also important to keep in mind that the gaugino condensate we have used for inflation does not play any role for low-energy SUSY breaking, as it becomes vanishing at the final (global) minimum. Hence, the magnitud of W (and thus of N_c) during inflation is not constrained by any low-energy SUSY breaking requirement, which would demand $W = \mathcal{O}(1\ TeV) M_{\rm Pl}^2$.

In conclusion, we have shown how the slow-roll problem, which generically affects all the inflationary models based on supergravity when the vacuum is dominated by an F-term, is naturally solved in the context of SUGRA theories coming form superstrings. In particular, this occurs if the vacuum energy during inflation is dominated by a modulus (T) F-term and the role of the inflaton is played by any field with the appropriate modular weight, namely $n = -3$. In this case, the inflaton mass is exactly vanishing at tree-level, and gets non-dangerous contributions at higher radiative levels. Inflation takes place in a hibrid-like manner. We have explicitely shown that it is easy to construct models implementing this idea successfully. In particular we have presented a model based on the vacuum energy

created by a gaugino condensate in the hidden sector, coupled to the inflaton field in a suitable way.

Acknowledgements

This research was supported in part by the CICYT (contract AEN95-0195) and the European Union (contract CHRX-CT92-0004) (JAC), and by the US Department of Energy under grant DE-FG03-91ER40662 Task C (GG).

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